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Superstring Phenomenology Present and Future Perspective¹

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Abstract. The objective of superstring phenomenology is to develop the models and methodology needed to connect quantitatively between Planck scale physics and electroweak scale experimental data. I review the present status of this endeavor with a focus on the three generation free fermionic models.

1. Introduction

Superstring phenomenology aims at achieving two goals. The first is to reproduce the physics of the Standard Model. The second is to identify possible experimental signature of superstring unification which may provide further evidence for its validity. A model which satisfies all of the experimental constraints, is likely to be more than an accident. Such a model, or class of models, will then serve as the laboratory for the search for exotic predictions of superstring unification, and as a laboratory in which we can address the important question of how the string vacuum is selected. I would like to remark that the goal of superstring phenomenology should not necessarily be viewed as “to derive the correct string vacuum”, but to develop the models and methodology needed to connect quantitatively between physical phenomena, which in the framework of unification

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is believed to originate at the Planck scale, and experimental data at the electroweak scale. One should not forget that physics is primarily an experimental science, and as current data seem to support the existence of unification, this endeavor is required.

The first question that we must ask is why and whether the pursuit of superstring theory is justified. After all one should not overlook the fact that currently it is assumed that gravity and the gauge interactions become compatible only near the Planck scale, which is some seventeen orders of magnitude above the scale where experiments are performed. Thus, superstring unification, based on input from current experiments, necessarily ignores the possibility that the real physics above the electroweak scale is entirely different from what we may have naively envisioned.

Ever since the Pati–Salam [1] and the Georgi–Glashow [2] seminal papers, we know that the gauge groups and matter content of the Standard–Model fit into representations of larger unifying gauge groups. This is an amazing coincidence that correlates at least eighteen experimental observable, if one just counts the quantum numbers of the Standard Model matter states. The match is so striking that it is hard to believe that nature is merely playing a cruel mirage. Proton lifetime constraints impose that such higher unification must occur at a scale which is at least of the order 10^{16-17}GeV , just one or two orders of magnitude below the Planck scale. Furthermore, low energy measurements of the Standard Model gauge coupling are in qualitative agreement with the idea of unification [3, 4], while precision measurements of the Standard Model parameters at the electroweak scale impose severe constraints on various extensions of the Standard Model. Notable among these extensions is weak scale supersymmetry, an important component in the program of unification, which continues to be in agreement with the experimental data. Thus, it is reasonable to hold the view that current low energy data support the notion of the big desert scenario and unification.

While unification and supersymmetry are very attractive concepts they fall short of being satisfactory. Many fundamental questions still cannot be addressed within their domain. The first type of questions is that supersymmetry and unification cannot answer why the Standard Model has the structure that it has. Namely, what is the mechanism that selects the gauge group and the matter content of the Standard Model. Furthermore, although some of the low energy couplings are manifestation of unified couplings at the high energy scale, unification and supersymmetry, in general, cannot explain the observed fermion mass spectrum. Most important, traditional point quantum field theories and gravity are incompatible. Thus, although gauge unification occurs one or two orders of magnitude below the Planck scale, gravity is still not incorporated. This is a crucial drawback of conventional grand unified theories as gravitational effects start to play

a deterministic role near the unification scale.

In the context of unified theories the solution to these problems will come from some fundamental theory at the Planck scale where gravity and quantum gauge field theories become of comparable strength. This elusive theory will hopefully tell us how the parameters of the Standard Model obtain their values. To date, superstring theory is the most developed Planck scale theory. It is believed to provide a framework for the consistent unification of gravity and the gauge interactions. As such it is currently the best probe that we have to explore physics at the Planck scale. It therefore makes sense to try to see if it is possible to connect superstring theory with the Standard Model. This is the subject of superstring phenomenology.

2. Different approaches to superstring phenomenology

There are currently two complementary approaches to superstring phenomenology. In the first, which can be regarded as a view from above, the general strategy is to first try to understand what is the nonperturbative formulation of string theory. The hope is that the string vacuum will then be uniquely fixed and the low energy predictions unambiguously determined. The second, which can be regarded as the view from below, asserts that we must use low energy data to single out phenomenologically interesting superstring vacua. Such string models will then be instrumental to understand the dynamics which select the string vacuum. These two approaches are in a sense complementary and progress is likely to be made by pursuing both approaches in parallel. Another view which I heard from Henry Tye suggests that just as our earth and our sun are not at the center of the universe, perhaps there does not exist a unique string vacuum and in effect we happen to live in a typical phenomenological string vacuum. If this is the correct view then while it still makes sense to look for the fundamental physical principles which underly string theory, or in more general terms, which underly Planck scale physics, string dynamics may not be able to reveal to us what is the superstring vacuum of our world.

I think that the discussion in the previous paragraph demonstrates a different point in regard to the present status of superstring phenomenology. The point is that the inclination toward one approach or another arises from the age old question of “what is the right question ?” That is what is the question whose answer will lead to real progress in understanding the objective structure of the fundamental particles and forces. Some people believe that the right question is “What is string theory ?”, and hence what is the fundamental non perturbative formulation of string theory. Others may believe that the important question is “how does unification occur ?” and regard, for example, the MSSM gauge coupling unification

as strong support for addressing this question. Yet another possibly important question is “how is supersymmetry broken?”. Another point of view is expressed concisely in a recent paper of Pati [5]. In this view the important questions are “how can the proton be stable?”, and “how can the neutrino masses be so small?”. The point is that if we assume that unification exist, then generically unification gives rise to many sources for proton decay and may also generate large neutrino masses. Experimentally the proton lifetime and the neutrino masses are suppressed to many orders of magnitude. Thus, these two experimental observations provide the most severe constraints on any model of unification. To satisfy those constraints simultaneously in a satisfactorily robust way are in this view the important questions to address.

In non supersymmetric grand unification proton decay is induced by the gauge bosons of the grand unifying gauge group. The experimental constraints then restrict that the unification scale, and the masses of the grand unifying gauge bosons, have to be above roughly 10^{16}GeV . The naive models then cannot be simultaneously consistent with the gauge coupling unification and proton stability. Complicating the spectrum may assist by pushing the unification to a higher scale. This is precisely what supersymmetry does. Supersymmetry, however, introduces many new sources of proton decay through dimension four and five operators. These new channels, in a generic supersymmetric model, will cause the proton to decay extremely fast. Thus, very severe constraints must be imposed on these new supersymmetric operators. In field theoretic supersymmetric theories this is achieved by imposing R -parity, which forbids dimension four operators, while a non trivial doublet-triplet splitting mechanism is needed to evade proton decay from dimension five operators. Thus, in field theoretic supersymmetric theories it may be possible to circumvent the problem. However, the type of symmetries that must be imposed are not expected to survive once we try to embed supersymmetry into a theory that unifies gravity with the gauge interactions. In gravity unified theories, to paraphrase one of the speakers in this conference, “everything which is not explicitly forbidden is allowed”. Similarly the field theoretic doublet-triplet splitting mechanism will be difficult, if not impossible, to implement in a theory that unifies gravity with the gauge interactions. Gravity unified theories would generically also give rise to operators which induce large neutrino masses, far above the experimental limits [6].

A search for a suitable resolution of these issues is likely to single out the phenomenologically interesting models of unification. Superstring models offer appealing solutions to these problems. In regard to proton stability superstring theory allows the existence of unification without the need for an enlarged grand unifying gauge group in the effective field theory level. In this respect we may have a grand unifying gauge group, which is broken

at the string level rather than in the effective field theory level. This is an appealing situation because in that case the doublet–triplet splitting problem is resolved by a superstringy doublet–triplet splitting mechanism [7]. In this superstringy mechanism color triplets that mediate rapid proton decay are simply projected out from the massless spectrum by the generalized GSO projections. Thus, dimension five operators can only be induced by heavier string modes. In specific models it is also possible to show that also heavy string modes do not induce such operators and in these models proton decay from dimension five operators is entirely forbidden. As for dimension four operators, again in the generic situation they are abundant and are difficult to avoid. The point is that higher order nonrenormalizable terms may induce dimension four operators that are not sufficiently suppressed. Their absence can only be insured by a gauge symmetry or a local discrete symmetry that will guarantee that the nonrenormalizable terms do not appear to a sufficient high order.

3. Superstring constructions

There are several possible ways to try to construct realistic superstring models. One possibility is to construct superstring models with an intermediate GUT gauge group, like $SU(5)$, $SO(10)$, E_6 , etc, which is broken to the Standard Model gauge group at an intermediate energy scale [8]. This option imposes that the intermediate non–Abelian gauge group is realized as a higher level affine Lie algebra because level one models do not contain adjoint Higgs representations in the massless spectrum.

The second possibility is to construct superstring models with semi–simple GUTs, like $SU(3)^3$ [9], $SU(5) \times U(1)$ [11, 12] or $SO(6) \times SO(4)$ [13]. In these models the extended non–Abelian symmetry is broken in the effective field theory at an intermediate energy scale. This type of models can be realized with $k \geq 1$ affine Lie algebra.

The last possibility is to construct superstring models in which the non–Abelian factors of the Standard–Model gauge group are obtained directly at the string levels [14, 15, 16, 17, 18, 19]. These models again can be realized with $k \geq 1$ affine Lie algebra. I propose that proton decay constraints suggest that this last possibility is the correct choice.

The general goal is therefore to construct superstring models that are as realistic as possible. A realistic model of unification must satisfy a large number of constraints a few of which are listed below.

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1. Gauge group $\longrightarrow SU(3) \times SU(2) \times U(1)_Y$
 2. Contains three generations
 3. Proton stable $(\tau_P > 10^{30+} \text{ years})$

- 4. N=1 supersymmetry (or N=0)
- 5. Contains Higgs doublets \oplus potentially realistic Yukawa couplings
- 6. Agreement with $\sin^2 \theta_W$ and α_s at M_Z (+ other observables).
- 7. Light left-handed neutrinos
 - 8. $SU(2) \times U(1)$ breaking
 - 9. SUSY breaking
 - 10. No flavor changing neutral currents
 - 11. No strong CP violation
 - 12. Exist family mixing and weak CP violation
- 13. + ...
- 14. + **GRAVITY**

The question then is whether it is possible to construct a model which satisfies all of those criteria, or possibly a class of models which can accommodate most of these constraints. To date the most developed theory that can consistently unify gravity with the gauge interactions is string theory. While alternatives may exist, it makes sense at this stage to try to use string theory to construct a model which satisfies the above requirements. Even if eventually string theory turns out not to be the fundamental theory of nature, a model which satisfies all of above constraints is likely to arise as an effective model from the true fundamental theory.

There are five known superstring theories in ten dimensions. The type I, type II A&B and the $E_8 \times E_8$ and $SO(32)$ heterotic strings. All these 10 dimensional string theories and including the 11 dimensional supergravity are believed to be special limits of one underlying theory. This is the understanding which emerges from the various string dualities that were uncovered in the last couple of years. The formulation of the underlying theory is still unclear and the hope is that it will improve our understanding of the mechanism which selects the string vacuum in four dimensions. In terms of trying to connect superstring theory with experimental physics the best that we can do at the moment is to continue to study promising four dimensional vacua and keep an open eye on progress in the formal understanding of superstring theory.

The exploration of realistic superstring vacua proceeds by studying compactification of the heterotic string from ten to four dimensions. There is a large number of possibilities. The first type of semi-realistic superstring vacua that were constructed are compactification on Calabi-Yau manifolds which give rise to an E_6 observable gauge group which is broken further by the Hosotani flux breaking mechanism to $SU(3)^3$ [9]. This gauge group is then broken to the Standard Model gauge group in the effective field

theory level. This type of geometrical compactifications are not exact conformal field theories and correspond at special points to conformal theories which have $N = 2$ world-sheet supersymmetry in the left and right moving sectors. Similar geometrical compactifications which have only $(2,0)$ world-sheet supersymmetry have also been studied and can lead to compactifications with $SO(10)$ and $SU(5)$ observable gauge group [10]. The analysis of this type of compactification is complicated due to the fact that they do not correspond to free world-sheet theories. Therefore, it is rather difficult to try to calculate the parameters of the Standard Model in these compactifications. On the other hand they can provide a more sophisticated mathematical window to the underlying quantum geometry and indeed much of the effort in this direction has focused on the formal development of the underlying quantum geometry [20].

The next class of superstring vacua that have been explored in detail are the orbifold models [21]. In these models one starts with a compactification of the heterotic string on a flat torus, using the Narain prescription [22]. This type of compactification uses free world-sheet bosons. The Narain lattice is then moded out by some discrete symmetries which are the orbifold twisting. The most detailed study of this type of models are the Z_3 orbifold [14] which give rise to three generation models with $SU(3) \times SU(2) \times U(1)^n$ gauge group. One caveat of this class of models is that the weak-hypercharge does not have the standard $SO(10)$ embedding. Thus, the nice features of $SO(10)$ unification are lost. This fact has a crucial implication that the normalization of the weak hypercharge relative to the non-Abelian currents is larger than $5/3$, the standard $SO(10)$ normalization. This results generically in disagreement with the observed low energy values for $\sin^2 \theta_W(M_Z)$ and $\alpha_s(M_Z)$.

Another special type of string compactifications that has been studied in detail are the free fermionic models. The simplest examples correspond to $Z_2 \times Z_2$ orbifolds at special points in the compactification space. These models give rise to the most realistic superstring models constructed to date. There are several key features of these models that suggest that if string theory is relevant in nature, then the true string vacuum is in the vicinity of these models. First is the fact that the free fermionic models are formulated at a highly symmetric point in the string compactification space. The second is that in the $Z_2 \times Z_2$ orbifold models at the free fermionic point the emergence of three generations is correlated with the underlying structure of the orbifold model [23]. Each of the Standard Model generations is obtained from one of the twisted sectors and carries horizontal charges under one of the orthogonal planes of the $Z_2 \times Z_2$ orbifold model. Naively one may view the existence of three generations in nature as arising because we are dividing the six dimensional compactified space into factors of two. Thus, these models may explain the existence of three generation

in nature in terms of the underlying geometry. The last important point is that in the free fermionic models the weak hypercharge has the standard $SO(10)$ embedding. Consequently these models can be in agreement with the observed low energy values for $\sin^2 \theta_W(M)$ and $\alpha_s(M_Z)$. This class of models and their phenomenology are the focus of this talk.

4. Free fermionic models

In the free fermionic formulation [24] for the left-movers (world-sheet supersymmetric) one has the usual space-time fields X^μ , ψ^μ , ($\mu = 0, 1, 2, 3$), and in addition the following eighteen real free fermion fields: χ^I, y^I, ω^I ($I = 1, \dots, 6$), transforming as the adjoint representation of $SU(2)^6$. The supercurrent is given in terms of these fields as follows

$$T_F(z) = \psi^\mu \partial_z X_\mu + \sum_{i=1}^6 \chi^i y^i \omega^i.$$

For the right movers we have \bar{X}^μ and 44 real free fermion fields: $\bar{\phi}^a$, $a = 1, \dots, 44$. A model in this construction is defined by a set of basis vectors of boundary conditions for all world-sheet fermions, which are constrained by the string consistency requirements and completely determine the vacuum structure of the model. The physical spectrum is obtained by applying the generalized GSO projections. The low energy effective field theory is obtained by S-matrix elements between external states. The Yukawa couplings and higher order nonrenormalizable terms in the superpotential are obtained by calculating correlators between vertex operators. For a correlator to be non vanishing all the symmetries of the model must be conserved. Thus, the boundary condition vectors determine the phenomenology of the models.

The first five vectors in the basis of the models that I discuss here consist of the NAHE³ set. The gauge group after the NAHE set is $SO(10) \times E_8 \times SO(6)^3$ with $N = 1$ space-time supersymmetry, and 48 spinorial 16 of $SO(10)$, sixteen from each sector b_1 , b_2 and b_3 . The NAHE set divides the internal world-sheet fermions in the following way: $\bar{\phi}^{1, \dots, 8}$ generate the hidden E_8 gauge group, $\bar{\psi}^{1, \dots, 5}$ generate the $SO(10)$ gauge group, and $\{\bar{y}^{3, \dots, 6}, \bar{\eta}^1\}$, $\{\bar{y}^1, \bar{y}^2, \bar{\omega}^5, \bar{\omega}^6, \bar{\eta}^2\}$, $\{\bar{\omega}^{1, \dots, 4}, \bar{\eta}^3\}$ generate the three horizontal $SO(6)^3$ symmetries. The left-moving $\{y, \omega\}$ states are divided to $\{y^{3, \dots, 6}\}$, $\{y^1, y^2, \omega^5, \omega^6\}$, $\{\omega^{1, \dots, 4}\}$ and $\chi^{12}, \chi^{34}, \chi^{56}$ generate the left-moving $N = 2$ world-sheet supersymmetry.

³ This set was first constructed by Nanopoulos, Antoniadis, Hagelin and Ellis (NAHE) in the construction of the flipped $SU(5)$. *nahe*=pretty, in Hebrew.

The internal fermionic states $\{y, \omega|\bar{y}, \bar{\omega}\}$ correspond to the six left-moving and six right-moving compactified dimensions in a geometric formulation. This correspondence is illustrated by adding the vector X to the NAHE set, with periodic boundary conditions for the set $(\bar{\psi}^1, \dots, \bar{\psi}^5, \bar{\eta}^{1,2,3})$ and antiperiodic boundary conditions for all other world-sheet fermions. This boundary condition vector extends the gauge symmetry to $E_6 \times U(1)^2 \times E_8 \times SO(4)^3$ with $N = 1$ supersymmetry and twenty-four chiral 27 of E_6 . The same model is generated in the orbifold language [21] by moding out an $SO(12)$ lattice by a $Z_2 \times Z_2$ discrete symmetry with standard embedding. In the construction of the standard-like models beyond the NAHE set, the assignment of boundary conditions to the set of internal fermions $\{y, \omega|\bar{y}, \bar{\omega}\}$ determines many of the properties of the low energy spectrum, such as the number of generations, the presence of Higgs doublets, Yukawa couplings, etc.

In the realistic free fermionic models the boundary condition vector X is replaced by the vector 2γ in which $\{\bar{\psi}^1, \dots, \bar{\psi}^5, \bar{\eta}^1, \bar{\eta}^2, \bar{\eta}^3, \bar{\phi}^1, \dots, \bar{\phi}^4\}$ are periodic and the remaining left- and right-moving fermionic states are antiperiodic. The set $\{1, S, 2\gamma, \xi_2\}$ generates a model with $N = 4$ space-time supersymmetry and $SO(12) \times SO(16) \times SO(16)$ gauge group. The b_1 and b_2 twists are applied to reduce $N = 4$ to $N = 1$ space-time supersymmetry. The gauge group is broken to $SO(4)^3 \times U(1)^3 \times SO(10) \times E_8$. The $U(1)$ combination $U(1) = U(1)_1 + U(1)_2 + U(1)_3$ has a non-vanishing trace and the trace of the two orthogonal combinations vanishes. The number of generations is still 24, eight from each sector b_1 , b_2 and b_3 . The chiral generations are now 16 of $SO(10)$ from the sectors b_j ($j = 1, 2, 3$). The $10 + 1$ and the E_6 singlets from the sectors $b_j + X$ are replaced by vectorial 16 of the hidden $SO(16)$ gauge group from the sectors $b_j + 2\gamma$. As I will show below the structure of the sector $b_j + 2\gamma$ with respect to the sectors b_j plays an important role in the texture of fermion mass matrices. The anomalous $U(1)$ in this models is seen to arise due to the breaking of $(2, 2) \rightarrow (2, 0)$ world-sheet supersymmetry. This anomalous $U(1)$ generates a Fayet-Iliopoulos D-term [25] which breaks supersymmetry at the Planck scale. Supersymmetry is restored by assigning non vanishing VEVs to a set of Standard Model singlets in the massless string spectrum along flat F and D directions.

The standard-like models are constructed by adding three additional vectors to the NAHE set [16, 17, 18, 19]. The $SO(10)$ symmetry is broken in two stages, first to $SO(6) \times SO(4)$ and next to $SU(3) \times SU(2) \times U(1)^2$. At the same time the number of generations is reduced to three generations one from each sector b_1 , b_2 and b_3 . The flavor $SO(6)^3$ symmetries are broken to factors of $U(1)_{R_j}$ and $U(1)_{R_{j+3}}$ ($j = 1, 2, 3$). For every right-moving gauged $U(1)$ symmetry there is a corresponding left-moving global $U(1)$ symmetry $U(1)_{L_j}$ and $U(1)_{L_{j+3}}$. Finally, each generation is charged

with respect to two Ising model operators which are obtained by pairing a left-moving real fermion with a right-moving real fermion. The Higgs spectrum consists of three pairs from the Neveu-Schwarz sector, h_j and \bar{h}_j , with charges under the horizontal symmetries and one or two additional pairs from the sector $b_1 + b_2 + \alpha + \beta$. Analysis of the Higgs mass matrix suggests that at low energies only one pair of Higgs multiplets remains light. The pair which remains light, together with the flavor symmetries which are broken near the string scale, then determines the hierarchical structure of the fermion mass matrices.

5. Fermion masses

One of the important achievements of the free fermionic models is the successful prediction of the top quark mass and the explanation of the heaviness of the top quark mass relative to the masses of the lighter quarks and leptons. The cubic level Yukawa couplings for the quarks and leptons are determined by the boundary conditions in the vector γ according to the following rule [19]

$$\Delta_j = |\gamma(U(1)_{L_{j+3}}) - \gamma(U(1)_{R_{j+3}})| = 0, 1 \quad (j = 1, 2, 3) \quad (1)$$

$$\Delta_j = 0 \rightarrow d_j Q_j h_j + e_j L_j h_j; \quad (2)$$

$$\Delta_j = 1 \rightarrow u_j Q_j \bar{h}_j + N_j L_j \bar{h}_j, \quad (3)$$

where $\gamma(U(1)_{R_{j+3}})$, $\gamma(U(1)_{L_{j+3}})$ are the boundary conditions of the world-sheet fermionic currents that generate the $U(1)_{R_{j+3}}$, $U(1)_{L_{j+3}}$ symmetries. In models with $\Delta_{1,2,3} = 1$ only $+\frac{2}{3}$ charged type quarks get a cubic level Yukawa coupling.

Such models therefore suggest an explanation for the top quark mass hierarchy relative to the lighter quarks and leptons. At the cubic level only the top quark gets a mass term and the mass terms for the lighter quarks and leptons are obtained from nonrenormalizable terms. To study this scenario we have to examine the nonrenormalizable contributions to the doublet Higgs mass matrix and to the fermion mass matrices [26, 27].

At the cubic level there are two pairs of electroweak doublets. At the nonrenormalizable level one additional pair receives a superheavy mass and one pair remains light to give masses to the fermions at the electroweak scale. Requiring F-flatness imposes that the light Higgs representations are \bar{h}_1 or \bar{h}_2 and h_{45} .

The nonrenormalizable fermion mass terms of order N are of the form $cgf_i f_j h \phi^{N-3}$ or $cgf_i f_j \bar{h} \phi^{N-3}$, where c is a calculable coefficient, g is the gauge coupling at the unification scale, f_i , f_j are the fermions from the sectors b_1 , b_2 and b_3 , h and \bar{h} are the light Higgs doublets, and ϕ^{N-3} is

a string of Standard Model singlets that get a VEV and produce a suppression factor $(\langle\phi\rangle/M)^{N-3}$ relative to the cubic level terms. Several scales contribute to the generalized VEVs. The leading one is the scale of VEVs that are used to cancel the “anomalous” $U(1)$ D-term equation. The next scale is generated by Hidden sector condensates. Finally, there is a scale which is related to the breaking of $U(1)_{Z'}$ which is embedded in $SO(10)$ and is orthogonal to the weak hypercharge.

At the cubic level only the top quark gets a non vanishing mass term. Therefore only the top quark mass is characterized by the electroweak scale. The remaining quarks and leptons obtain their mass terms from nonrenormalizable terms. The cubic and nonrenormalizable terms in the superpotential are obtained by calculating correlators between the vertex operators. The top quark Yukawa coupling is generically given by, $g\sqrt{2}$, where g is the gauge coupling at the unification scale. In the model of Ref. [18], bottom quark and tau lepton mass terms are obtained at the quartic order,

$$W_4 = \{d_{L_1}^c Q_1 h'_{45} \Phi_1 + e_{L_1}^c L_1 h'_{45} \Phi_1 + d_{L_2}^c Q_2 h'_{45} \bar{\Phi}_2 + e_{L_2}^c L_2 h'_{45} \bar{\Phi}_2\}.$$

The VEVs of Φ are obtained from the cancelation of the anomalous $U(1)$ D-term equation. The coefficient of the quartic order mass terms were calculated by calculating the quartic order correlators and the one dimensional integral was evaluated numerically. Thus after inserting the VEV of $\bar{\Phi}_2$ the effective bottom quark and tau lepton Yukawa couplings are given by [18],

$$\lambda_b = \lambda_\tau = 0.35g^3.$$

They are suppressed relative to the top Yukawa by

$$\frac{\lambda_b}{\lambda_t} = \frac{0.35g^3}{g\sqrt{2}} \sim \frac{1}{8}.$$

To evaluate the top quark mass, the three Yukawa couplings are run to the low energy scale by using the MSSM RGEs. The bottom mass is then used to calculate $\tan\beta$ and the top quark mass is found to be [18],

$$m_t \sim 175 - 180 GeV.$$

The fact that the top Yukawa is found near a fixed point suggests that this is in fact a good prediction of the superstring standard-like models. By varying $\lambda_t \sim 0.5 - 1.5$ at the unification scale, it is found that λ_t is always $O(1)$ at the electroweak scale.

The analysis of the fermion masses of the lighter quarks and leptons proceeds by analyzing higher order nonrenormalizable terms. An analysis of fermion mass terms up to order $N = 8$ revealed the general texture of fermion mass matrices in these models [26]. The sectors b_1 and b_2 produce

the two heavy generations. Their mass terms are suppressed by singlet VEVs that are used in the cancelation of the anomalous $U(1)$ D-term equation. The sector b_3 produces the lightest generation. The diagonal mass terms for the states from b_3 can only be generated by VEVs that break $U(1)_{Z'}$. This is due to the horizontal $U(1)$ charges and because the Higgs pair h_3 and \bar{h}_3 necessarily gets a Planck scale mass [26]. The suppression of the lightest generation mass terms is seen to be a result of the structure of the vectors α and β with respect to the sectors b_1 , b_2 and b_3 . The mixing between the generations is obtained from exchange of states from the sectors $b_j + 2\gamma$. The general texture of the fermion mass matrices in the superstring standard-like models is of the following form,

$$M_U \sim \begin{pmatrix} \epsilon, a, b \\ \tilde{a}, A, c \\ \tilde{b}, \tilde{c}, \lambda_t \end{pmatrix}; \quad M_D \sim \begin{pmatrix} \epsilon, d, e \\ \tilde{d}, B, f \\ \tilde{e}, \tilde{f}, C \end{pmatrix}; \quad M_E \sim \begin{pmatrix} \epsilon, g, h \\ \tilde{g}, D, i \\ \tilde{h}, \tilde{i}, E \end{pmatrix},$$

where $\epsilon \sim (\Lambda_{Z'}/M)^2$. The diagonal terms in capital letters represent leading terms that are suppressed by singlet VEVs, and $\lambda_t = O(1)$. The mixing terms are generated by hidden sector states from the sectors $b_j + 2\gamma$ and are represented by small letters. They are proportional to $\langle TT \rangle / M^2$.

In Ref. [27] it was shown that if the states from the sectors $b_j + 2\gamma$ obtain VEVs in the application of the DSW mechanism, then a Cabibbo angle of the correct order of magnitude can be obtained in the superstring standard-like models. For one specific choice of singlet VEVs that solves the cubic level F and D constraints the down mass matrix M_D is given by

$$M_d \sim \begin{pmatrix} \epsilon & \frac{V_2 \bar{V}_3 \Phi_{45}}{M^3} & 0 \\ \frac{V_2 \bar{V}_3 \Phi_{45} \xi_1}{M^4} & \frac{\bar{\Phi}_2^- \xi_1}{M^2} & 0 \\ 0 & 0 & \frac{\Phi_1^+ \xi_2}{M^2} \end{pmatrix} v_2, \quad (4)$$

where $v_2 = \langle h_{45} \rangle$ and we have used $\frac{1}{2}g\sqrt{2\alpha'} = \sqrt{8\pi}/M_{Pl}$, to define $M \equiv M_{Pl}/2\sqrt{8\pi} \approx 1.2 \times 10^{18} GeV$ [28]. The undetermined VEVs of $\bar{\Phi}_{13}$ and ξ_2 are used to fix m_b and m_s such that $\langle \xi_1 \rangle \sim M$. We also take $\tan\beta = v_1/v_2 \sim 1.5$. Substituting the values of the VEVs above and diagonalizing M_D by a bi-unitary transformation we obtain the Cabibbo mixing matrix

$$|V| \sim \begin{pmatrix} 0.98 & 0.2 & 0 \\ 0.2 & 0.98 & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (5)$$

Since the running from the scale M down to the weak scale does not affect the Cabibbo angle by much [29], we conclude that realistic mixing of the correct order of magnitude can be obtained in this scenario. The analysis was extended to show that reasonable values for the entire CKM matrix

parameters can be obtained for appropriate flat F and D solutions. For one specific solution the up and down quark mass matrices take the form

$$M_u \sim \begin{pmatrix} \epsilon & \frac{V_3 \bar{V}_2 \Phi_{45} \bar{\Phi}_3^+}{M^4} & 0 \\ \frac{V_3 \bar{V}_2 \Phi_{45} \bar{\Phi}_2^+}{M^4} & \frac{\bar{\Phi}_i^- \bar{\Phi}_i^+}{M^2} & \frac{V_1 \bar{V}_2 \Phi_{45} \bar{\Phi}_2^+}{M^4} \\ 0 & \frac{V_1 \bar{V}_2 \Phi_{45} \bar{\Phi}_1^+}{M^4} & 1 \end{pmatrix} v_1, \quad (6)$$

and

$$M_d \sim \begin{pmatrix} \epsilon & \frac{V_3 \bar{V}_2 \Phi_{45}}{M^3} & 0 \\ \frac{V_3 \bar{V}_2 \Phi_{45} \xi_1}{M^4} & \frac{\bar{\Phi}_2^- \xi_1}{M^2} & \frac{V_1 \bar{V}_2 \Phi_{45} \xi_i}{M^4} \\ 0 & \frac{V_1 \bar{V}_2 \Phi_{45} \xi_i}{M^4} & \frac{\Phi_1^+ \xi_2}{M^2} \end{pmatrix} v_2, \quad (7)$$

with v_1, v_2 and M as before. Substituting the VEVs and diagonalizing M_u and M_d by a bi-unitary transformation, we obtain a qualitatively realistic mixing matrix. The texture and hierarchy of the mass terms in Eqs. (6,7) arise due to the set of singlet VEVs in the solution to the F and D -flatness constraints. The zeroes in the 13 and 31 entries of the mass matrices are protected to all orders of nonrenormalizable terms. To obtain a non-vanishing contribution to these entries either V_1 and \bar{V}_3 or V_3 and \bar{V}_1 must obtain a VEV simultaneously. Thus, there is a residual horizontal symmetry that protects these vanishing terms. The 11 entry in the mass matrices, e.g. the diagonal mass terms for the lightest generation states, can only be obtained from VEVs that break $U(1)_{Z'}$ [26]. We assume that $U(1)_{Z'}$ is broken at an intermediate energy scale that is suppressed relative to the scale of scalar VEVs [26]. In Ref. [30] we showed that $U(1)_{Z'}$ is broken by hidden sector matter condensates at $\Lambda_{Z'} \leq 10^{14} GeV$. Consequently, we have taken $\epsilon \leq (\Lambda_{Z'}/M)^2 \sim 10^{-8}$.

Texture zeroes in the fermion mass matrices are obtained if the VEVs of some states from the sectors $b_j + 2\gamma$ vanish. These texture zeroes are protected by the symmetries of the string models to all order of nonrenormalizable terms [27]. For example in the above mass matrices the 13 and 31 vanish because $\{V_1, V_3\}$ get a VEV but \bar{V}_1 and \bar{V}_3 do not. Other textures are possible for other choices of VEVs for the states from the sectors $b_j + 2\gamma$.

6. Gauge coupling unification

An important issue in superstring phenomenology is that of gauge coupling unification. Superstring theory predicts that the gravitational and gauge couplings are unified and satisfy the relation,

$$\frac{8\pi G_N}{\alpha'} = g_i^2 k_i \equiv g_{\text{string}}^2 \quad (8)$$

where G_N is the gravitational coupling, α' is the Regge slope and k_i are the Kac-Moody level of the group factor G_i . Thus, superstring theory predicts that all gauge couplings are unified without the need for a grand unified gauge group. This is an important property of superstring theory. For while superstring theory still predicts unification of the couplings, and for this reason many of the successful relations of grand unified theories can be retained, many of the problems associated with GUTs, like too rapid proton decay etc., can be resolved. The string unification scale [31] is of the order,

$$M_{\text{string}} \approx 5 \times g_{\text{string}} \times 10^{17} \text{GeV} \quad (9)$$

If we naively assume that the spectrum below the string unification scale consists solely of the MSSM spectrum, *i.e.* three generations plus two Higgs doublets and the weak hypercharge normalization is $k_Y = 5/3$, then the predicted values for $\sin^2 \theta_W(M_Z)$ and $\alpha_s(M_Z)$ are in contradiction with the experimentally observed values. Alternatively we may identify the MSSM unification scale of the order $2 \times 10^{16} \text{GeV}$ which is roughly a factor of 20 below the string unification scale. This discrepancy is known as the string gauge coupling unification problem.

It would seem that in an extrapolation of the gauge couplings of fifteen orders of magnitude, there will be many possible resolutions to this problem. In this regard we must caution that one should view such fine structure of the theory with a grain of salt. While we believe that our method of extrapolation is consistent and safe, one should not forget that many assumptions are being made with regard to the nature of the physics in the desert between the weak and unification scales. With this remark in mind, we find that in fact the string coupling unification problem is not easily resolved. To understand why this is the case, it suffices to examine the extrapolation of α_s . We know that starting from the MSSM unification scale and extrapolating to the Z scale we get a value of order $\alpha_s \approx 0.12$, in agreement with the data. Now suppose that we extrapolate further from the weak scale to the bottom scale. The strong coupling at the bottom scale is then of order 0.22, which agrees with the data. Now the effect of starting the running at the string scale is exactly the same. Namely, instead of adding the additive interval between the Z and the bottom scales, we add it at the high scale between the MSSM and the string unification scales. Therefore, we find that $\alpha_s(M_Z) \approx 0.22$, which is roughly 100% off the experimentally observed value. We see that although the discrepancy between the string and MSSM unification scales is small, it translates into a large deviation in α_s . This is the reason why most of the sources that may affect the evolution of α_s cannot, in fact, resolve the problem.

In ref. [32] the string gauge coupling unification problem was examined in detail. From the gauge couplings renormalization group equations we obtain solutions for $\sin^2 \theta_W(M_Z)$ and $\alpha_s(M_Z)$, which include all the possi-

ble sources that may affect their values at M_Z . These include: modification of the weak hypercharge normalization [33], extended non-Abelian gauge group at an intermediate energy scale, threshold corrections due to the infinite tower of the heavy string modes, threshold corrections due to the light SUSY spectrum, the effect of additional matter thresholds at intermediate energy scales. The analysis also includes the two-loop gauge and Yukawa coupling corrections to the central values, as well as the effect of converting from the \overline{MS} renormalization scheme in which the parameters are measured to the \overline{DR} scheme in which the SUSY parameters are run.

The most complex part of the analysis is the calculation of the string threshold corrections. This analysis is made complicated by the fact that a typical realistic superstring model contains a few hundred-thousand sectors over which we have to integrate. In addition one needs to worry about careful removal of the infrared divergences and the integration over the modular domain. Although in simple unrealistic models it is possible to parameterize the string threshold corrections in terms of moduli fields [34], this is in general not possible in the realistic models which are by far more complicated. A priori we have no reason to expect that the string threshold corrections would be small. It is a viable possibility that the string thresholds would produce large corrections that may shift the string unification scale to the MSSM unification scale. The calculation of the string threshold corrections is similar to the calculation of the one-loop partition function. While in the partition function we count all states at each mass level, the string threshold corrections count all states at each mass level times their lattice charge. Technical details of the calculation are given in ref. [32].

The conclusion of ref. [32] is that of the perturbative corrections to the one-loop RGEs, only the existence of additional matter thresholds [35] at intermediate energy scale may result in agreement of the string unification scale with the low energy data. All other effects are either too small or in fact increase the value of $\alpha_s(M_Z)$ and therefore worsen the disagreement with the low energy data. The effect of intermediate color triplets is to sufficiently slow down the evolution of α_s , while all other effects are either too small or increase $\alpha_s(M_Z)$. This is a remarkable conclusion as it shows that perturbative string scale unification requires the existence of additional matter beyond the MSSM spectrum. String models generically give rise to additional matter beyond the MSSM spectrum, and some models in fact produce the additional representations with the required weak hypercharge assignment, needed for the consistency of string scale gauge coupling unification with the low energy data [36]. Another proposal [37] suggests that nonperturbative effects in the context of M-theory are responsible for shifting the string unification scale down to the MSSM unification scale.

7. Exotic matter

The existence of additional matter beyond the MSSM spectrum may have important phenomenological and cosmological implications [39]. The extra matter is often obtained in the realistic free fermionic models from sectors which arise due to the breaking of the $SO(10)$ symmetry to one of its subgroups. This matter can be classified according to the type of $SO(10)$ breaking in each sector. Sectors which break the $SO(10)$ symmetry to $SO(6) \times SO(4)$ or to $SU(5) \times U(1)$ give rise to states with fractional charge $\pm 1/2$ while sectors which break the $SO(10)$ symmetry directly to $SU(3) \times SU(2) \times U(1)^2$ give rise to states with the Standard Model charges but with fractional charges under $U(1)_{Z'}$ which is embedded in $SO(10)$. Thus, for example we have states from this type of sectors which are exotic leptoquarks [38] *i.e.*, they have baryon number $\pm 1/3$ but exotic lepton number $\pm 1/2$. Due to its exotic $U(1)_{Z'}$ charge, in some models, it can be shown that the superpotential terms of such a state with the Standard Model states vanish to all orders of nonrenormalizable terms. In this case the exotic states can interact with the Standard Model states only via the gauge interactions and cannot decay into them. In such a model therefore an exotic state will be stable and one has to check that its mass density does not over close the universe. These constraints were investigated in detail in ref. [39]. Several general remarks however are in order. Exotic states that do not have GUT origin are generic in superstring models. They arise due to the breaking of the non-Abelian gauge symmetries at the string rather than in the effective field theory level. Such states are therefore a generic signature of superstring compactification. Thus, they may lead to possible observable experimental signatures. For example, all the level one models predict the existence of fractionally charged states at least with Planck scale masses. Specific free fermionic models also predict the existence of Standard Model states which are exotic from the point of view of the underlying $SO(10)$. These may be color triplets, electroweak doublets or Standard Model singlets. Such states may provide an experimental signature of specific classes of superstring compactifications.

8. Supersymmetry breaking

Next let me turn to the question of supersymmetry breaking. All the models that I discussed so far have $N = 1$ supersymmetry at the Planck scale. Supersymmetry must be broken at some scale as it is not observed at the low scale.

we address the following question: Given a supersymmetric string vacuum at the Planck scale, is it possible to obtain hierarchical supersymmetry

breaking in the observable sector? A supersymmetric string vacuum is obtained by finding solutions to the cubic level F and D constraints. We take a gauge coupling in agreement with gauge coupling unification, thus taking a fixed value for the dilaton VEV. We then investigate the role of nonrenormalizable terms and strong hidden sector dynamics. The hidden sector contains two non-Abelian hidden gauge groups, $SU(5) \times SU(3)$, with matter in vector-like representations. The hidden $SU(3)$ group is broken near the Planck scale. We analyze the dynamics of the hidden $SU(5)$ group. The $SU(5)$ hidden matter mass matrix is given by

$$\mathcal{M} = \begin{pmatrix} 0 & C_1 & 0 \\ B_1 & A_2 & C_2 \\ 0 & C_3 & A_1 \end{pmatrix}, \quad (10)$$

where A, B, C arise from nonrenormalizable terms of orders $N = 5, 8, 7$ respectively. A specific solution was found in ref. [40] and taking generically $\langle \phi \rangle \sim gM/4\pi \sim M/10$ yielded $A_i \sim 10^{15} \text{ GeV}$, $B_i \sim 10^{12} \text{ GeV}$, and $C_i \sim 10^{13} \text{ GeV}$. From Eqs. (10) we observe that to insure a nonsingular hidden matter mass matrix, we must require $C_1 \neq 0$ and $B_1 \neq 0$. This imposes $\bar{V}_3 \neq 0$ and $V_2 \neq 0$. Thus, the nonvanishing VEVs that generate the Cabibbo mixing also guarantee the stability of the supersymmetric vacuum. The gaugino and matter condensates are given by the well known expressions for supersymmetric $SU(N)$ with matter in $N + \bar{N}$ representations [41],

$$\begin{aligned} \frac{1}{32\pi^2} \langle \lambda\lambda \rangle &= \Lambda^3 \left(\det \frac{\mathcal{M}}{\Lambda} \right)^{1/N}, \\ \Pi_{ij} = \langle \bar{T}_i T_j \rangle &= \frac{1}{32\pi^2} \langle \lambda\lambda \rangle \mathcal{M}_{ij}^{-1} \end{aligned} \quad (11)$$

where $\langle \lambda\lambda \rangle$, \mathcal{M} and Λ are the hidden gaugino condensate, the hidden matter mass matrix and the $SU(5)$ condensation scale, respectively. Modular invariant generalization of Eqs. (11) for the string case were derived in Ref. [42]. The nonrenormalizable terms can be put in modular invariant form by following the procedure outlined in Ref. [43]. Approximating the Dedekind η function by $\eta(\hat{T}) \approx e^{-\pi\hat{T}/12}(1 - e^{-2\pi\hat{T}})$, we verified that the calculation using the modular invariant expression from Ref. [42] (with $\langle \hat{T} \rangle \approx M$) differ from the results using Eq. (11), by at most an order of magnitude. The hidden $SU(5)$ matter mass matrix is nonsingular for specific F and D flat solutions. In Ref. [40] a specific cubic level F and D flat solution was found. The gravitino mass due to the gaugino and matter condensates was estimated to be of the order $1 - 10 \text{ TeV}$. The new aspect of our scenario for supersymmetry breaking is that supersymmetry is broken due to the interplay of the scale generated by the anomalous $U(1)$ and the inclusion of nonrenormalizable terms. Hidden sector strong dynamics at an

intermediate scale may then be responsible for generating the hierarchy in the usual way. Using field theoretic toy models it was demonstrated that supersymmetry can indeed be broken in this fashion [44].

9. R -parity violation

Turning to a different issue, recently the H1 and ZEUS collaborations reported an excess in $e^+P \rightarrow e^+\text{jet}$ events [45]. A possible interpretation of this excess is the existence of supersymmetry with R -parity violating couplings. This requires that the lepton number violating couplings are of order one while proton lifetime restrictions impose that the baryon number violating couplings are suppressed by at least 22 orders of magnitude. Therefore we seek a mechanism which allows the lepton number violating couplings and forbids the baryon number violating couplings.

The model of ref. [46] provides an example how such a mechanism can be realized in superstring theory. This model gives rise to additional space-time vector bosons from the basis vectors $\{\alpha, \beta, \gamma\}$ which extend the NAHE set. These additional vector bosons enhance a combination of the universal and flavor $U(1)$ symmetries to a custodial $SU(2)$ symmetry. Only the Standard Model leptons transform under this $SU(2)$ symmetry while the quarks are singlets. As a result nonrenormalizable terms of the form $QLDN$ are allowed by gauge invariance while terms of the form $UDDN$ and $QQQL$ are forbidden to all orders of nonrenormalizable terms. Thus, the VEV of the right handed neutrino in this model induces the desired lepton number violating coupling while the baryon number violating couplings are forbidden to all orders of nonrenormalizable terms.

10. Beyond superstring

Next I briefly discuss superstring phenomenology in the context of the recently discovered superstring dualities. Superstring theory gives rise to several duality symmetries [47]. The simplest example which illustrates how these dualities arise is the example of compactification on a circle. When we compactify on a circle the fact that the wave-function is single valued imposes that the momenta on the compactified dimension are quantized. This produces the string momentum modes. In string theory in addition we can wrap the string on the compactified dimension which produces the winding modes. There is a duality symmetry which interchanges momentum and winding modes and at the same time $R \leftrightarrow 1/2R$. This simple one dimensional example extends to higher dimensions. In two dimensions we find three types of duality symmetries. The first which is geometrical in

nature, usually referred to as U -duality, the second which exchanges momentum and winding modes, typically referred to as T -duality and finally a symmetry which exchanges $T \leftrightarrow U$ and is known as mirror symmetry.

The dualities mentioned above are perturbative and exhibit themselves in the exchange of the spectrum and the superpotential. Thus they can be checked order by order in perturbation theory. In the last few years we have witnessed a significant progress in understanding duality symmetries which are nonperturbative, *i.e.* they exchange weak with the strong coupling. The starting point in this program is the Seiberg-Witten solution of $N = 2$ supersymmetric pure $SU(2)$ [48]. In the supersymmetric theory the gauge coupling is extended to a complex parameter $\tau = \theta/2\pi + i4\pi/g^2$ where θ is the axial coupling and g is the field strength coupling. The strong-weak duality extends to the full $SL(2, Z)$ duality of the parameter τ . In the Seiberg-Witten solution the exact vacuum structure of the theory is parameterized in terms of a genus one Riemann surface.

In string theory we have roughly and very naively a similar situation. The gauge coupling is fixed by the VEV of the dilaton field. The dilaton field, combined with the space-time components of the antisymmetric tensor field forms a modular parameter. In M-theory [50, 51] this complex field is identified with the moduli field of a new dimension and hence the $SL(2, Z)$ symmetry of this moduli field translates into a duality which exchanges strong and weak coupling [49]. In the last couple of years a large number of qualitative tests have been performed which confirm this basic picture. The hope that understanding nonperturbative string phenomena would reveal how a specific vacuum is selected did not materialize so far. Although the underlying formulation of M-theory is still unclear, preliminary attempts have been made to extract some phenomenological information related to it [52].

The study of nonperturbative dualities indicates a new structure underlying superstring theory and quantum space-time. Underlying superstring and M-theory there should exist a unifying theory which produces the five known ten dimensional strings and 11 dimension supergravity as special limits. Currently the popular interpretation is the Matrix model construction of M-theory [53]. Excellent reviews on these attempts exist in the literature [54]. Here I briefly review my work with Matone which may (or may not) be related to this interpretation of quantum space-time.

With Matone, inspired by some of the mathematical structure underlying the recently proposed nonperturbative dualities, we proposed a duality between the space-coordinate and the wave function in quantum mechanics [55, 56]. We suggested that in quantum mechanics the space coordinate should be interpreted as statistical variable of a thermodynamical theory underlying the quantum space-time. Carrying this interpretation a step further means that the space coordinate in quantum mechanics should be

replaced by a sort of density matrix. This interpretation is similar in spirit to the Matrix model approach to M-theory. The expected outcome is that following this interpretation of the space-coordinate in quantum mechanics a string-like structure would emerge. The reason is that the proposed $\{x-\psi\}$ duality suggests that second quantization of ψ implies an expansion of x which is reminiscent of the expansion of x in terms of the string modes. However, at present this is still very speculative and not well founded.

Another view of the recently discovered dualities is that different vacua that are classically distinct, are in fact equivalent quantum mechanically. Equivalence here means that there exist some coordinate transformation that takes one vacuum to another. Recently with Matone we formulated an equivalence principle in quantum mechanics [57]. Starting from the classical Hamilton-Jacobi equation for the reduced action we showed that equivalence under diffeomorphism requires modification of the classical equation which is the quantum Hamilton-Jacobi equation. This quantum equivalence principle then states that all the quantum mechanical systems with different potentials are equivalent under a general coordinate transformation. The relation of this equivalence principle to the recently discovered nonperturbative dualities is yet to be uncovered.

11. Conclusions and outlook

In summary, string theory provides a window to Planck scale physics and to the study of the unification of the gauge interactions with gravity. To bring this exploration from mere speculations into contact with experimental physics we have to construct phenomenological superstring models. The realistic free fermionic models achieved remarkable success in describing the real world: existence of three generations with standard $SO(10)$ embedding, a superstring solution to the doublet-triplet splitting problem, correct top quark mass prediction, potentially realistic fermion mass textures, agreement with $\alpha_s(M_Z)$ and $\sin^2 \theta_W(M_Z)$, etc. These successes of the realistic free fermionic models may be accidental. However, taking the view that that is not the case, the fact that free fermionic models are formulated at a highly symmetric point in the moduli space and the underlying $Z_2 \times Z_2$ orbifold compactification may be the origin of their realistic nature. Finally, perturbative string scale unification is possible provided that additional color triplets and electroweak doublets exist at intermediate energy scales. Such additional states may give rise to experimentally accessible signature which may provide the smoking gun of superstring unification.

As an outlook to the future of the phenomenology of superstring unification it should be stressed that the important problem in high energy physics is the nature of the electroweak symmetry breaking mechanism.

This question must be answered by our experimental colleagues and we eagerly await their resolution of the issue. With this understanding in mind, and with current experimental data, supersymmetry and unification are the most promising theoretical proposals. If indeed this is the avenue chosen by nature then most of the properties of the Standard Model parameters are determined by some fundamental Planck scale theory. Therefore, we must develop the technology needed to study Planck scale physics. The realistic free fermionic models are designed to achieve precisely this goal, being the most realistic superstring model constructed to date. At this time we have only begun to scratch the surface of this class of models and much further work is required both in order to develop the calculational tools needed to confront string models with the experimental data and to better understand the properties of this class of models. Understanding the correspondence with other formulations will provide further insight into their realistic properties, while experimental findings of exotic particles may provide further support for their validity. Finally, it is hoped that these explorations will assist in understanding if and how a particular string vacuum is selected and the nature of quantum space-time.

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